

Sequestering in String Theory

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We study sequestering, a prerequisite for flavor-blind supersymmetry breaking in several high-scale mediation mechanisms, in compactifications of type IIB string theory. We find that although sequestering is typically absent in unwarped backgrounds, strongly warped compactifications do readily sequester. The AdS/CFT dual description in terms of conformal sequestering plays an important role in our analysis, and we establish how sequestering works both on the gravity side and on the gauge theory side. We pay special attention to subtle compactification effects that can disrupt sequestering. Our result is a step toward realizing an appealing pattern of soft terms in a KKLT compactification.

Contents

| | |
|--|----|
| 1. Introduction | 1 |
| 1.1. A milder criterion for sequestering | 3 |
| 1.2. Outline | 5 |
| 2. Sequestering in unwarped backgrounds | 5 |
| 2.1. Restrictions on D3-brane couplings | 5 |
| 2.2. Non-sequestered couplings via compactification moduli | 7 |
| 2.3. Sequestering in the case of “localized” Kähler moduli | 11 |
| 2.4. Complex structure moduli | 12 |
| 2.5. Quantum corrections to sequestering | 13 |
| 3. Warped sequestering and AdS/CFT | 14 |
| 3.1. The background | 15 |
| 3.2. Conformal sequestering | 17 |
| 4. Gravity-side sequestering | 18 |
| 4.1. D3-brane couplings in a non-compact warped throat | 19 |
| 4.2. SUSY breaking and sequestering in a non-compact warped throat | 20 |
| 4.3. Compactification effects | 22 |
| 5. Classes of warped sequestering | 23 |
| 6. Conclusion | 25 |

1. Introduction

Several of the most promising ideas about high-scale transmission of supersymmetry (SUSY) breaking to the supersymmetric standard model (SM) involve the physics of extra dimensions. These ideas, which include anomaly mediation [1,2], gaugino mediation [3], and mirage mediation [4], each require the strong suppression of standard gravity-mediation effects, which would otherwise naively dominate. For this to happen robustly, it is important that the SUSY-breaking hidden sector be “sequestered” away from the visible SM fields in the extra-dimensional space.

Concretely, the troubling gravity-mediation term is really a direct cross-coupling in the effective Lagrangian written in curved superspace, between the hidden-sector fields X bearing the dominant F-term, and the visible-sector fields Q^i , of the form

$$\int d^4\theta \, X^\dagger X \, \frac{c_{ij}}{M_P^2} (Q^i)^\dagger Q^j . \quad (1.1)$$

For order-one and flavor-violating c_{ij} , phenomenologically dangerous flavor-violating squark and slepton mass terms would result. This is a version of the “supersymmetric flavor problem”. We say that X and Q have been sequestered from one another if $c_{ij} \ll 1$.

It is important to note [1] that the cross-couplings in curved superspace should be identified with cross-couplings within (1.1) in

$$f = -3M_P^2 \exp\left(-\frac{K}{3M_P^2}\right) \quad (1.2)$$

rather than within the effective four-dimensional supergravity Kähler potential, K , itself.

As originally envisioned, sequestering could be justified by locality in extra dimensions of space [1]. That is, if SUSY were broken on a brane separated from the SM brane by a distance d in the extra dimensions, cross-couplings in the four-dimensional effective theory should only arise from integrating out the exchange of massive bulk modes. If these modes all had sufficiently large masses $m \gg 1/d$, these bulk exchanges would be exponentially screened by Yukawa suppression, leading to $c_{ij} \sim e^{-m \cdot d} \ll 1$. In this circumstance, other effects, *e.g.* the anomaly mediated contribution $\sim \frac{1}{16\pi^2} g^2 \frac{|F|^2}{M_P^2}$, could provide dominant, and flavor-blind, soft masses.¹ An important subtlety is that at least some bulk modes should get masses $m \lesssim 1/d$ from compactification itself or from moduli stabilization, and therefore their exchanges require careful analysis. A thorough study was made within a minimal five-dimensional effective field theory with a single extra dimension, confirming that sequestering was indeed robust in that context [6].

Nevertheless, it has been argued that in string theory, with its much richer structure and several extra dimensions, it is very difficult to realize sequestering [7,8]. In particular, these works made it clear that spatial separation alone does *not* suffice to ensure sequestering in string constructions. However, these studies were not exhaustive and whether sequestering could be realized in a complete theory of quantum gravity remained an open and important question.

A different approach to sequestering was initiated in [9]. It was shown there, again within minimal five-dimensional effective field theory, that sequestering was consistent with strongly warped compactifications and their attendant geometric hierarchies. This suggested, via the AdS/CFT correspondence [10,11], that there should exist a dual four-dimensional mechanism for sequestering. In “conformal sequestering” [12], instead of a weakly coupled hidden sector, we imagine a hidden sector that is strongly coupled over some range of scales, and is close to a conformal fixed point over the range from Λ_{UV} to Λ_{IR} . (An alternate approach of conformal sequestering, where it is the visible sector that is

¹ For the particular case of anomaly mediation, additional mechanisms are required to avoid tachyonic sleptons, with concrete suggestions appearing in *e.g.* [1,5].

strongly coupled and conformal over a range of scales, was proposed earlier in [13].) At the scale Λ_{IR} , conformal invariance breaks, and soon after this, spontaneous supersymmetry breaking occurs. Some operator \mathcal{O} in the hidden sector CFT provides the leading coupling

$$\int d^4\theta \mathcal{O} c_{ij}(Q^i)^\dagger Q^j \quad (1.3)$$

of SUSY breaking to the visible squarks and sleptons. But now if \mathcal{O} has scaling dimension $2+\Delta$, one finds an additional suppression of the naive coupling (1.1) by a factor of $\left(\frac{\Lambda_{IR}}{\Lambda_{UV}}\right)^\Delta$. For a typical intermediate-scale scenario, this suppresses the dangerous cross-couplings sufficiently for anomaly mediation (or one of the other extra-dimensional mechanisms) to provide safely flavor-universal visible masses. Conformal sequestering has been further developed in [14,15,16].

In this paper we demonstrate that the mechanism of “warped sequestering” can be realized in string theory, in the form leading to anomaly-mediated supersymmetry breaking in the visible sector. In essence, we will describe how warped compactifications of string theory, involving warped throats similar to the canonical Klebanov-Strassler example [17], compactified as in [18], can naturally yield sequestered SUSY breaking. We will make considerable use of the AdS/CFT dual grammar of conformal sequestering, although the explicit construction will be on the gravity side of the duality. We also show that sequestering is possible in unwarped compactifications, but typically absent.

Our goal here is to establish that sequestering can be achieved fairly robustly in a theory of quantum gravity. We will not discuss the construction of fully realistic models. This would require both inclusion of a detailed visible sector, and incorporation of one of the various model-building mechanisms to *e.g.* fix the tachyonic sleptons of anomaly mediation.

We will restrict ourselves to type IIB string theory constructions, but our final conclusions will clearly be more general. More precisely, since we will always consider non-vanishing brane separations to be much larger than the string length, we will work in terms of type IIB effective supergravity coupled to branes.

1.1. A milder criterion for sequestering

While sequestering originally referred to suppression of *all* terms of the form (1.1), we will make use of a milder criterion following from a simple observation of [16]. This states that we can tolerate an unsuppressed bilinear hidden superfield appearing in (1.1) if

it has as its vector component a conserved Noether current for a hidden sector IR (non-R) symmetry, because then the hidden bilinear is guaranteed to have vanishing SUSY-breaking VEV. The symmetry is only required to be good symmetry of the hidden dynamics close to the SUSY breaking scale, *even if it is ultimately spontaneously broken by hidden VEVs*.

We discuss here a simple example of the power of this observation, and will have occasion to return to this example in the course of the paper. It is provided by considering a distribution of D3-branes in a simple toroidal orientifold T^6/Z_2 compactification of the six extra dimensions, preserving $\mathcal{N} = 4$ supersymmetry. Splitting the branes into two well-separated stacks, with chiral matter (in $\mathcal{N} = 1$ language) labeled Q and X , provides a toy model of visible and hidden sectors. While this X hidden sector does not break SUSY, it does allow us to study sequestering itself. This example was previously invoked as a case in which string theory does not sequester [7], as follows. The Kähler potential is readily determined by the high degree of supersymmetry and turns out to be

$$f = -3 \left[(S + S^\dagger) \det \left(T_{IJ} + T_{JI}^\dagger - \text{tr } Q_I Q_J^\dagger - \text{tr } X_I X_J^\dagger \right) \right]^{1/3}, \quad (1.4)$$

where $I, J = 1, 2, 3$ label the $\mathcal{N} = 1$ chiral superfields in an $\mathcal{N} = 4$ multiplet and the traces are over the different branes in each stack. We assume that the moduli S, T_{IJ} are supersymmetrically stabilized by some contributions to their superpotential (while this would not happen in the $\mathcal{N} = 4$ theory, it can happen in close relatives that inherit the relevant pieces of f). At low energies they can therefore be set to their SUSY-preserving VEVs. It is clear that expanding f in powers of brane fields will result in cross-couplings of the form (1.1).

But f does satisfy the milder version of sequestering, as we now check. Expanding (1.4) in powers of Q, X (setting the moduli to their VEVs), we get

$$f = \text{const.} - \text{tr } \tilde{Q}_I^\dagger \tilde{Q}^I - \text{tr } \tilde{X}_I^\dagger \tilde{X}^I + \text{tr } \tilde{Q}_I^\dagger \tilde{Q}^I \left[\delta_I^J \text{tr } \tilde{X}_K^\dagger \tilde{X}^K - 3 \text{tr } \tilde{X}_I^\dagger \tilde{X}^J \right] + \dots, \quad (1.5)$$

where the canonical fields are given by

$$\tilde{Q}_I = \text{const.} (T + T^\dagger)_{IJ}^{-1/2} Q_J, \quad (1.6)$$

$$\tilde{X}_I = \text{const.} (T + T^\dagger)_{IJ}^{-1/2} X_J. \quad (1.7)$$

The hidden bilinear, traceless in the I, J indices, manifestly corresponds to currents of the $SU(3)$ subgroup of the $\mathcal{N} = 4$ R-symmetry. From the $\mathcal{N} = 1$ viewpoint, $SU(3)$ is a non-R flavor symmetry acting on chiral multiplets. While the toroidal compactification breaks this $SU(3)$ in the UV, it is an accidental symmetry of the hidden sector in the IR and therefore this example satisfies the milder criterion of [16]. (Later we will obtain a more general understanding of why this happened.)

1.2. Outline

The remainder of the paper is organized as follows. In §2, we begin by studying unwarped backgrounds with D-branes, ultimately specializing to D3-branes. We show how unwarped compactification in general spoils sequestering, although we point out exceptional cases where sequestering holds to a good approximation. In §3, we study branes in highly warped backgrounds, in particular specializing to the case of the Klebanov-Strassler conifold throat. We discuss the AdS/CFT dual description of conformal sequestering and show that, even upon compactification, sequestering must hold. In §4 we illustrate some aspects of this sequestering directly from computations on the gravity side of the duality. We illustrate the conjunction of both hidden sector SUSY breaking and warped sequestering by the example of an anti-D3-brane at the end of the throat [19]. In §5 we discuss general classes and properties of warped throats which appear to be promising for warped sequestering. We conclude in §6.

2. Sequestering in unwarped backgrounds

In this section we discuss the possibility of sequestering in compactifications of type IIB string theory with little or no warping. Although we will eventually be led to pursue sequestering in warped compactifications, the unwarped case provides an instructive introduction to the general problem. In §2.1 we study D-branes coupled to supergravity in the simplest context, ten non-compact dimensions. In particular, we explain why the fields on spatially-separated D3-branes might be expected to enjoy sequestering. Then, in §2.2, we expose the fallacy in this logic: compactification effects do generically spoil sequestering of D3-branes in unwarped compactifications, albeit with some notable exceptions. We then describe the effects on sequestering from Kähler moduli corresponding to blowup modes (§2.3), from complex structure moduli (§2.4), and from quantum corrections to the Kähler potential (§2.5).

2.1. Restrictions on D3-brane couplings

We begin by recalling, from [7], a reason for fearing that two stacks of D-branes, say stack Q and stack X, will have dangerous cross-terms such as (1.1) in their four-dimensional effective Lagrangian. In string theory, D-branes source various ten-dimensional fields, such as the metric and dilaton, and thereby affect their surroundings. By affecting the background in this way, each stack can perceive the other. Specifically, denote the metric on

the compactification manifold by $g_{mn}(y)$, and the axio-dilaton as $\tau(y)$. Then backreaction of stack Q on g_{mn} and τ , which will implicitly make g and τ into functions of the scalar modes of Q , will translate directly into couplings between Q and X when one computes the probe action for the X branes in the same background.

As a concrete example, we can consider the toy model on p.40 of the second reference of [7]. Imagine that the X-stack is the hidden sector, and is composed of a Dp' brane. Let the Q-stack, where the visible sector resides, be the probe brane. The Dp' brane background metric and dilaton are given by

$$ds^2 = h(r)^{-1/2} dx_{||}^2 + h(r)^{1/2} dx_{\perp}^2, \quad e^{-2\phi} = h(r)^{(p'-3)/2} \quad (2.1)$$

where the harmonic function $h(r)$ is given by

$$h(r) = 1 + g_s \left(\frac{\sqrt{\alpha'}}{r} \right)^{7-p'}. \quad (2.2)$$

The action of brane Q in this background is given by expanding the DBI action

$$S_Q = -T_{Dp} \int d^{p+1}x \, e^{-\phi} \sqrt{\det(\gamma_{\mu\nu} + \dots)} \quad (2.3)$$

where $\gamma_{\mu\nu}$ is the induced world-volume metric, and the omitted terms involve worldvolume fluxes that are unimportant at present. Expanding in the background (2.1) yields

$$S_Q \propto \int d^{p+1}x \, \sqrt{\det(\tilde{\gamma}_{\mu\nu})} \, h(r)^{(p'-p)/4-1} \left(1 + \frac{1}{2} h(r) \partial_{\mu} \phi^i \partial^{\mu} \phi_i + \dots \right), \quad (2.4)$$

where $\tilde{\gamma}_{\mu\nu}$ is the worldvolume metric induced from the *unwarped* background metric.

Interestingly, when the solution takes the form (2.1) for $p = p'$, the leading effect of the backreaction of the X-branes on the Q-branes does not alter the kinetic terms.² This is the question of interest for sequestering, since the dangerous term (1.1) would imply $|X|^2 |\partial Q|^2$ scalar cross-couplings (among other terms).

This example is at best a toy model for the full conformally Calabi-Yau solutions with nontrivial fluxes of the NSNS three-form H_3 and the RR three-form F_3 and five-form F_5 . The general supergravity solution in the presence of these fluxes is discussed in some detail

² The appearance of a potential in (2.3) in this case is spurious: it is canceled by the Chern-Simons terms in the action.

in [18] (building on the earlier works [20,21,22]). The main points, at leading order in the g_s and α' expansions, are the following:

- The metric takes the form

$$ds^2 = e^{2A(y)} \eta_{\mu\nu} dx^\mu dx^\nu + e^{-2A(y)} \tilde{g}_{mn}(y) dy^m dy^n \quad (2.5)$$

where the warp factor $A(y)$ depends on the D3-brane positions and the D7-brane positions, as well as on the background flux.

- The dilaton takes the form $\tau = \tau(y)$, and is determined self-consistently in terms of \tilde{g}_{mn} and the positions of any D7-branes. There is *no* explicit dependence on the D3-brane positions.
- The Einstein equations determining the internal metric \tilde{g}_{mn} depend on the dilaton gradients and the localized D7-brane actions, but again have *no* explicit dependence on the D3-brane positions.

It follows from these facts that the cancellation that prevents cross-couplings of the form (1.1) between D3-brane fields from different stacks in the simple ansatz (2.1) will continue to hold for the probe analysis of kinetic terms of a given D3-brane stack in the full supergravity solutions of [18]. While the D3-branes see any background D7-branes via their kinetic terms, they do not see each other at this order. In fact, the D3-branes also enjoy a vanishing potential, even though generic fluxes do break supersymmetry.

We note that this reasoning does not work for inter-D7-brane couplings in the same class of type IIB solutions. The fortuitous cancellation of metric and dilaton factors that occurs for D3-branes probing other D3-branes in the background (2.1), does not occur for general Dp-brane pairs. It is a happy fact of life for the D3-branes that their positions only enter the supergravity data via factors of the warp factor e^A , and these factors cancel in the leading terms of the probe action for a D3-brane.

2.2. Non-sequestered couplings via compactification moduli

The fact that separated stacks of D3-branes are sequestered in non-compact ten-dimensional spacetime might lead us to expect the same feature after compactification, at least to leading order in α' and g_s . However, we now study the four-dimensional supergravity theory after compactification and show that it typically leads to complete breakdown of sequestering in backgrounds of the form of §2.1. While we limit ourselves in this section to Calabi-Yau orientifolds with D3-branes, incorporating D7-branes would not be difficult. Some of the results here were described (or are implicit) in [18,7,23,24,25,26,27,28].

The loophole in our previous discussion is that compactification results in moduli whose parametrization in terms of $\mathcal{N} = 1$ chiral superfields depends on the locations of any D3-branes. Couplings of these moduli to the D3-branes can thereby mediate non-sequestered cross-couplings between distant branes. We begin by explaining this in the case of a single breathing mode for the entire compact space.

We now use the ansatz

$$ds^2 = h^{-1/2} e^{-6u} g_{\mu\nu} dx^\mu dx^\nu + h^{1/2} e^{2u} \tilde{g}_{ij} dy^i dy^j. \quad (2.6)$$

Here $h(y) = e^{-4A(y)}$ is the usual warp factor, e^{6u} is the breathing mode of the compactification manifold M , and \tilde{g} is the unwarped metric, with fixed fiducial volume \tilde{V}_6 . The powers of e^u are chosen so that $g_{\mu\nu}$ is the spacetime metric in four-dimensional Einstein frame.

We will now show (following [28]) that the breathing mode depends on the D3-brane positions. We will then demonstrate that this leads to contact terms between separated D3-branes. The perturbation to the warp factor, δh , sourced by a single D3-brane at x obeys

$$-\nabla_y^2 \delta h(x; y) = 2\kappa_{10}^2 T_3 \left(\frac{\delta^{(6)}(x - y)}{\sqrt{g}} - \rho(y) \right), \quad (2.7)$$

where $\rho(y)$ is the background density of D3-brane charge, arising from fluxes and possibly from other branes. Here y is the coordinate on the internal space, and x is the position of the D3-brane on this space. The distinction is important: the four-dimensional action arising from dimensional reduction does not depend on y , but x is a scalar field in this action — x is a coordinate on the D3-brane moduli space. (This moduli space has the same metric as the physical compact space, at least at leading order.) Finally, the D3-brane tension is

$$T_3 = \frac{1}{(2\pi)^3 g_s \alpha'^2}, \quad (2.8)$$

and the ten-dimensional gravitational coupling is

$$\kappa_{10}^2 = \frac{1}{2} (2\pi)^7 g_s^2 \alpha'^4. \quad (2.9)$$

We solve (2.7) by first solving

$$-\nabla_{y'}^2 \Phi(y; y') = -\nabla_y^2 \Phi(y; y') = \frac{\delta^{(6)}(y - y')}{\sqrt{g}} - \frac{1}{V_6} \quad (2.10)$$

where $V_6 \equiv \int d^6 y \sqrt{g} = e^{6u} \tilde{V}_6$. The solution to (2.7) is then

$$\delta h(x; y) = 2\kappa_{10}^2 T_3 \left[\Phi(x; y) - \int d^6 y' \sqrt{g} \Phi(y; y') \rho(y') \right]. \quad (2.11)$$

An important consequence [28] is that

$$-\nabla_x^2 \delta h(x; y) = 2\kappa_{10}^2 T_3 \left[\frac{\delta^{(6)}(x - y)}{\sqrt{g(x)}} - \frac{1}{V_6} \right], \quad (2.12)$$

for any background charge distribution $\rho(y)$. Away from the D3-brane we may drop the $\delta^{(6)}$ term in (2.12), which leads to

$$\tilde{g}^{i\bar{j}} \partial_i \partial_{\bar{j}} \delta h = \left(\frac{\kappa_{10}^2 T_3}{\tilde{V}_6} \right) e^{-4u} \equiv \frac{T_3}{M_P^2} e^{-4u}, \quad (2.13)$$

where the power of u comes from converting V_6 to \tilde{V}_6 and g^{ij} to \tilde{g}^{ij} , and we have assumed that M is Kähler to simplify the Laplacian to the form shown.

The right hand side of (2.13) is a constant independent of y , but $\tilde{g}^{i\bar{j}}$ depends on y . This is consistent only if δh obeys

$$\partial_i \partial_{\bar{j}} \delta h = \frac{T_3}{3M_P^2} e^{-4u} \tilde{g}_{i\bar{j}}. \quad (2.14)$$

We conclude that (to leading order in the gravitational coupling [28])

$$\delta h = \frac{T_3}{3M_P^2} e^{-4u} k(x, \bar{x}) + \dots, \quad (2.15)$$

where k is the Kähler potential for the metric \tilde{g} on M , and the dots denote terms annihilated by the Laplacian, which are not constrained by the above argument.

Next, we use $h + \delta h$ to identify the appropriate holomorphic coordinate T on the Kähler moduli space. As explained in detail in [27,28], the appropriate definition that is consistent with (2.15) is

$$T + \bar{T} \equiv \int_{\Sigma} d^4 y \sqrt{G} = \int_{\Sigma} d^4 y \sqrt{\tilde{g}} h e^{4u}, \quad (2.16)$$

where $G_{ij} \equiv h^{1/2} \tilde{g}_{ij} e^{2u}$ is the full metric on the compact space. Here Σ is an appropriate four-cycle, and because we have one Kähler modulus, Σ is effectively unique. We may normalize the fiducial metric \tilde{g} so that

$$\int_{\Sigma} d^4 y \sqrt{\tilde{g}} = 1. \quad (2.17)$$

Next, we choose for simplicity a case without background warping, so that $h = 1 + \delta h$ gets its only nontrivial contribution from the D3-branes in question. Combining these ingredients with (2.15), we find

$$T + \bar{T} = e^{4u} + \gamma k(x, \bar{x}), \quad (2.18)$$

where

$$\gamma \equiv \frac{T_3}{3M_P^2}, \quad (2.19)$$

which makes manifest the D3-brane-dependence of the holomorphic parametrization of the breathing mode. (As discussed in [25], invariance of this expression under Kähler transformations of k requires that T also transforms.) It follows (using the standard relation between the Kähler potential for volume moduli and the compactification volume) that the Kähler potential is

$$K = -3 \log(T + \bar{T} - \gamma k(x, \bar{x})). \quad (2.20)$$

By repeating the above argument for an additional D3-brane located at x' , we find

$$K = -3 \log(T + \bar{T} - \gamma k(x, \bar{x}) - \gamma k(x', \bar{x}')). \quad (2.21)$$

Notice that this result is sequestered: the coefficient of 3 implies that the different brane stacks and T appear additively in f . This feature will, however, turn out to be a special case, as we now demonstrate by extending the above considerations to a configuration with multiple Kähler moduli and two stacks of D3-branes. The condition (2.12) is unchanged, as it depends only on the overall breathing mode, not the relative sizes of cycles. It follows that (2.15) is also unchanged. Next, if $\{\Sigma_\alpha\}$ is a basis of four-cycles, we can identify the proper holomorphic coordinates T_α ,

$$T_\alpha + \bar{T}_\alpha \equiv \int_{\Sigma_\alpha} \sqrt{\tilde{g}} e^{4u} (1 + \delta h(x, x')). \quad (2.22)$$

The Kähler-invariant coordinates are then

$$U_\alpha = T_\alpha + \bar{T}_\alpha - \gamma k_\alpha(x, \bar{x}) - \gamma k_\alpha(x', \bar{x}') \quad (2.23)$$

where the transformation properties of $k_\alpha \sim k \int_{\Sigma_\alpha} \sqrt{\tilde{g}}$ are dictated by [29]. The Kähler potential is then

$$K = -2 \log(V_6) = -2 \log(d^{\alpha\beta\gamma} U_\alpha^{1/2} U_\beta^{1/2} U_\gamma^{1/2}) \quad (2.24)$$

where $d^{\alpha\beta\gamma}$ is determined by the intersection form of the four-cycles $\{\Sigma_\alpha\}$.

From (2.23), (2.21) it is easy to see that non-sequestered terms between the D3-branes are indeed generated in compactifications with multiple Kähler moduli, but not in the single-modulus case.

As we discussed in §1 and will revisit in §4.3, the $\mathcal{N} = 4$ toroidal orientifold has multiple Kähler moduli but provides an interesting exception to the rule: this configuration exhibits sequestering in the subtler sense of [16]. However, general multi-moduli compactifications are not expected to be safe in this sense.

Our argument for (2.23) (*cf.* [28]) was slightly indirect, involving the background charge in (2.7) and the associated transformation of the holomorphic coordinates T_α under Kähler transformations. However, an earlier and far more explicit derivation of this same expression appears in [23]. Nevertheless, our method is instructive: we have seen that the mechanism by which D3-branes couple to each other in f is by mixing with the holomorphic Kähler moduli T_α , and, crucially, this mixing is a result of the *compactness* of the space, as it follows from the Gauss’s law constraint (2.12).

2.3. Sequestering in the case of “localized” Kähler moduli

The explicit results of [23] suggest that sequestering is spoiled in typical models with multiple Kähler moduli. Here we point out the existence of a class of multi-moduli models in which approximate sequestering can arise.

Consider the following example with two Kähler moduli. The Calabi-Yau hypersurface in $\mathbb{WP}_{1,1,1,6,9}^4$ has (as derived in *e.g.* [30])

$$K \propto -2 \log \left(U_1^{3/2} - U_2^{3/2} \right) . \quad (2.25)$$

Very roughly, one should think of U_1 as the overall volume of the bulk Calabi-Yau manifold, and of U_2 as controlling the volume of a blow-up mode for a $\mathbb{C}^3/\mathbb{Z}_3$ singularity localized in the bulk. Resolving the singularity introduces a finite-volume \mathbb{P}^2 in the geometry, with a smooth neighborhood locally modeled on the total space of the bundle $\mathcal{O}(-3) \rightarrow \mathbb{P}^2$. It is implausible, at least for small values of the blow-up mode, that introducing this little “bump” on the Calabi-Yau space should destroy sequestering between separated D3-branes, which one would have expected to hold in the absence of this blow-up. How can we see this in detail?

As explained in [23], in particular in their expression (3.13), in the general multi-moduli case the D3-brane fields appearing in (2.23) are multiplied by the harmonic form

corresponding to the Kähler modulus U_α *evaluated at the position on the compact manifold* of the given D3-brane stack. Therefore, if the compact manifold is a large space with point-like singularities controlled by several additional Kähler moduli, for generic D3-brane positions, the corresponding U_α will *not* have significant dependence on the D3-brane positions. That is, the wavefunction of the Kähler mode has a negligible overlap with the D3-brane.

In the two-modulus case introduced above, one would then find a Kähler potential of the form

$$K = -2 \log \left(U_1^{3/2} - d \right) \quad (2.26)$$

where d is nearly *independent* of the D3-brane modes. This is still not of the sequestered form, but one can do a systematic expansion in powers of d/Vol , in which the leading term shows sequestering, and in which the corrections are parametrically small. This is in accord with the intuition that an arbitrarily small blow-up controlled by an additional Kähler modulus should not completely destroy the sequestering that is present in the single modulus case.

In the context of warped compactifications an argument along these lines has been given in [31], where it was argued that the wavefunction of the bulk Kähler moduli has negligible overlap with the small three-cycle at the tip of a Klebanov-Strassler throat, thereby resulting in sequestering. We will find that this argument is not quite justified, but sequestering does nevertheless hold in the more relaxed sense of [16], as discussed in the introduction.

2.4. Complex structure moduli

We have not yet discussed the role of complex structure moduli, but we now show that they do not alter sequestering. For instance, consider a compactification with a single Kähler modulus, which, as we showed in §2.2, enjoys sequestering. In detail, the Kähler potential (at leading order in α' and g_s) with complex structure moduli z^i , a Kähler modulus corresponding to the harmonic two-form $\omega_{m\bar{n}}$, axio-dilaton τ , and D3-brane moduli ϕ^m is given by [23]

$$\frac{K_{tot}}{M_P^2} = K_{CS}(z) - \log \left(-i(\tau - \bar{\tau}) \right) - 2 \log \left(\frac{1}{6} \mathcal{K}(T, z, \phi) \right) . \quad (2.27)$$

Here

$$K_{CS}(z) = -\log \left(-i \int_M \Omega \wedge \bar{\Omega} \right) , \quad (2.28)$$

with Ω the unique (up to scale) holomorphic three-form on the Calabi-Yau space, and

$$-2 \log \mathcal{K} = -3 \log \frac{2}{3} \left(T + \bar{T} + \frac{3i}{2\pi} \omega_{m\bar{n}} \text{Tr}(\phi^m \bar{\phi}^{\bar{n}}) + \frac{3}{8\pi} (\omega_{m\bar{n}} \bar{z}^i (\bar{\chi}_i)_{\bar{l}}^{\bar{n}} \text{Tr}(\phi^m \phi^{\bar{l}}) + h.c.) \right), \quad (2.29)$$

where χ_i is related to a complex structure deformation of the compact manifold in a way that is not important for our discussion.

We concluded that sequestering holds because when the D3-brane fields ϕ^m are split into two stacks, Q and X , the dependences on Q and X appear additively in

$$f = -3M_P^2 \exp(-K_{tot}/3M_P^2), \quad (2.30)$$

without cross-couplings between fields living on distinct D3-branes.

The reader might, however, wonder about the role of the complex structure moduli z^i , which evidently couple directly to the D3-brane fields ϕ in (2.29). In flux models, the z^i can be stabilized supersymmetrically by the fluxes at some high scale $\sim \alpha'/\sqrt{V_6}$. Does integrating out these fields then lead to dangerous cross-couplings of the D3-brane fields?

The answer is no. After appropriately rescaling the ϕ fields to obtain canonical normalization (by Taylor expanding the logarithm and rescaling by the appropriate power of the volume), the couplings we are concerned about look like

$$\int d^4\theta \ (\bar{z}\phi\phi + h.c.). \quad (2.31)$$

When we integrate out the z fields, it is easy to see that we will introduce contact terms between ϕ fields on different D3-branes, but these terms will be of higher dimension than (1.1). We can safely ignore these higher-derivative terms.

2.5. Quantum corrections to sequestering

Let us continue to focus on the model of a single Kähler modulus discussed in the previous subsection. The exact Kähler potential receives non-sequestered α' and g_s corrections to (2.29). The leading correction inducing communication between separated D3-branes was discussed in [32], based on explicit string loop calculations in certain toroidal orientifolds. The relevant operator takes the form

$$\Delta K = \frac{c}{T + \bar{T}} \left(\frac{q(\phi, \bar{\phi})}{\tau - \bar{\tau}} \right) \quad (2.32)$$

where q is some local function of the D3-brane fields. It is evident that through the function (2.30), this correction gives rise to a contact term between fields on different D3-branes, arising from

$$\Delta f \sim \frac{c g_s}{T + \overline{T}} k(\phi, \bar{\phi}) q(\phi, \bar{\phi}) , \quad (2.33)$$

with k the leading-order D3-brane Kähler potential. However, as expected for a string loop correction, this term is parametrically small at weak string coupling. (The power of $T + \overline{T}$ does not represent a suppression, as this is part of the four-dimensional Planck mass suppression in (1.1) which is by itself insufficient for sequestering.)

The constant c in (2.32),(2.33) is important. Loop corrections in higher-dimensional supergravity and string theory are suppressed as in naive dimensional analysis, and in fact the value of c inferred from orientifold calculations [32],

$$c = \frac{1}{128\pi^6} , \quad (2.34)$$

indicates that sequestering is not only a parametric fact in some such models, but can also be numerically effective, even for reasonably large values of g_s .

3. Warped sequestering and AdS/CFT

In the previous section we found that in compactifications with little warping, sequestering is in general possible only in models with a single Kähler modulus.³ The absence of sequestering, despite the extra-dimensional separation of visible and hidden brane stacks by some distance d , is ultimately accounted for by a coupling mediated by bulk KK modes, which we have integrated out in arriving at the four-dimensional effective supergravity. Since such non-sequestered contributions are not seen from purely five-dimensional analyses in which the fifth dimension separates the hidden and visible sectors [6,9], the KK modes in question must be due to the other extra dimensions, with some characteristic radius $\sim R$. When all compactification scales are comparable, the KK mediation is unsuppressed, as we found above except in special circumstances. However, if there is a geometric hierarchy $d \gg R$, KK mediation should be Yukawa-suppressed $\sim e^{-d/R}$, and we would expect sequestering to arise. Warped throats provide a class of suitably asymmetric

³ The approximate sequestering of §2.3 is present in the limit in which all the four-cycles associated to additional Kähler moduli are blown down, which may be thought of as a one-modulus limit.

compactifications, and are readily realizable in string theory. In this section, we demonstrate that configurations of D3-branes separated along a warped throat naturally enjoy sequestering. The AdS/CFT correspondence provides a valuable tool in establishing this result, and also in making the connection to conformal sequestering.

We begin by describing one particular warped throat background of interest, then establish sequestering from the perspective of the gauge theory dual to this warped throat. We return to illustrate the central elements from the gravity viewpoint. This realization of sequestering is powerful enough to accommodate complex and realistic visible sectors engineered in Calabi-Yau bulk portions of warped compactifications, not just the simple gauge theories on D3-branes that we will consider here. We will not, however, pursue visible-sector model building in this paper; for the state of the art, see *e.g.* [33,34].

3.1. The background

Our setting is a type IIB warped flux compactification with the metric (2.5). Such warped Calabi-Yau compactifications arise naturally in type IIB string theory, in the presence of three-form fluxes in the internal space (and a five-form flux that is determined by the warp factor).

A particularly simple and explicit example is due to Klebanov and Strassler [17]. We start with a conifold, a singular Calabi-Yau space of the form

$$\sum_{i=1}^4 z_i^2 = 0 \tag{3.1}$$

in \mathbb{C}^4 . Upon deformation, *i.e.* adding a small constant ε to the RHS of (3.1)), one finds two natural three-cycles in the geometry: a three-sphere A that shrinks in the singular limit $\varepsilon \rightarrow 0$, and its dual three-cycle B . When these cycles are threaded by three-form fluxes

$$\int_A F_3 = M, \quad \int_B H_3 = -K \tag{3.2}$$

with $M \gg 1$ and $K \gg 1$, one finds a simple supergravity solution with many interesting properties. The flux superpotential deforms (3.1) to an equation with an exponentially small constant on the RHS. Far from the associated smoothed tip, the solution takes the form of a warped metric (2.5) on the Calabi-Yau cone over the Einstein manifold $T^{1,1}$, with

$$h(r) = e^{-4A(r)} = \frac{27\pi}{4r^4} \alpha'^2 g_s M K + \dots, \tag{3.3}$$

where the dots denote known logarithmic corrections that are unimportant for our considerations.

Focusing on the AdS-like portion of the metric, this is just a spacetime of the familiar form

$$ds^2 = \frac{r^2}{R^2}(-dt^2 + dx_i^2) + \frac{R^2}{r^2}dr^2, \quad (3.4)$$

with r the radial direction of the throat, and

$$R^4 = \frac{27}{4}\pi g_s N \alpha'^2, \quad N \equiv KM. \quad (3.5)$$

The minimal redshift in this throat arises at the smoothed tip (where the metric, although well-known [17], departs from the form we have shown here), and is [18]

$$\frac{r_0}{R} = \exp\left(-\frac{2\pi K}{3g_s M}\right). \quad (3.6)$$

Further, we can glue this throat region into a compact Calabi-Yau around $r \approx r_{max}$ for some large r_{max} , as in [18].

The field theory dual of the Klebanov-Strassler solution is an $SU(N+M) \times SU(N)$ gauge theory with chiral multiplets $A^{1,2}$ and $B^{1,2}$ in the representations $(N+M, \bar{N})$ and $(\overline{N+M}, N)$, respectively. The theory has a superpotential

$$W = \epsilon_{ij}\epsilon_{kl}\text{Tr}A^i B^k A^j B^l \quad (3.7)$$

which is invariant under an $SU(2) \times SU(2)$ global symmetry, where one factor rotates the A_i and the other rotates the B_j . This symmetry is broken in the UV by the compactification, but it is an accidental symmetry in the IR of the CFT, or equivalently far down the throat.

If the fixed point were a free field theory, (3.7) would be irrelevant. Instead, we should interpret this theory as follows. For $M = 0$, this theory is a CFT, dual to the supergravity background created by N D3-branes at a conifold singularity [35]. The chiral fields A, B have conformal dimension $3/4$ at the nontrivial fixed point. The theory with finite $M \ll N$ should be viewed in a $1/N$ expansion about this interacting fixed point. This results, as argued convincingly in [17], in a “renormalization group cascade.” If $N = KM$, this cascade occurs over a range of energy scales with

$$\frac{\Lambda_{IR}}{\Lambda_{UV}} = \exp\left(-\frac{2\pi K}{3g_s M}\right), \quad (3.8)$$

corresponding precisely to (3.6).

3.2. Conformal sequestering

Let us imagine that we can modify the background described in §3.1 in such a way that in the IR, at the scale Λ_{IR} , SUSY is broken. Various ways of doing this in the dual gravity solution have been proposed in *e.g.* [19,36], and we will discuss these in §4. The main important fact will be that the leading operators generating soft terms are highly irrelevant. Whether or not the SUSY breaking spontaneously breaks the $SU(2) \times SU(2)$ symmetry of the Klebanov-Strassler sector (conifold throat), the SUSY breaking will be sequestered from a visible sector localized elsewhere on the Calabi-Yau, as we now explain.

We follow the discussion in [16]. Quite generally, the leading correction to sequestering will be of the form

$$\int d^4\theta \, \hat{\mathcal{O}} \, Q^\dagger Q \quad (3.9)$$

where $\hat{\mathcal{O}}$ is the hidden sector *non-chiral operator with the smallest conformal dimension*. In practice, a dimension $\Delta \gtrsim 3$ provides adequate sequestering. If the hidden CFT has a non-R symmetry, there will be dimension-two non-chiral operators in the same supermultiplet as the associated Noether currents [12,37], but as mentioned in the introduction, these operators will be harmless if the relevant CFT deformations (that enable SUSY breaking) also preserve this symmetry [16]. (The $U(1)_R$ itself is harmless, never mediating problematic interactions; it is dual to the bulk graviphoton, which is similarly harmless [6]). The Klebanov-Strassler field theory is a deformation of the Klebanov-Witten [35] CFT by a small number of field theory colors, and does not alter the $SU(2) \times SU(2)$ global symmetry, so we conclude that there is no danger from the associated non-chiral operators.

The leading potentially dangerous non-chiral operators are then determined by the properties of the conifold throat in the gravity dual. The spectrum of bulk KK modes in $T^{1,1}$ compactification of type IIB supergravity was determined in [38] and has been summarized in a very useful way in Appendix A of [39]. The main results for our purposes are the following:

- It might seem that the lowest-dimension non-chiral supermultiplets that can mediate SUSY breaking to the visible sector are the $SU(2) \times SU(2)$ singlets $|A|^2$ and $|B|^2$, which have naive scaling dimension two. It would then appear that these can multiply visible-sector fields and break sequestering completely. However, the true scaling dimensions of $|A|^2, |B|^2$ grow with the 't Hooft coupling in the supergravity limit $g_s N \gg 1$, *i.e.* these operators are dual to string states in the conifold throat. On the other hand, the $SU(2) \times SU(2)$ non-singlet bilinears in A, B are just the harmless Noether currents discussed above.

- In fact, the lowest-dimension non-chiral supermultiplet of operators invariant under $SU(2) \times SU(2)$ is given by

$$\mathcal{O}_8 = W_\alpha^2 \overline{W}_{\dot{\alpha}}^2, \quad (3.10)$$

where W_α is the standard gauge field strength superfield whose θ expansion begins with the gaugino λ_α . The lowest component of \mathcal{O}_8 has dimension $\Delta = 6$, while the highest component (relevant for transmission of scalar soft masses in (3.9)) has $\Delta = 8$. While these are the expected free-field dimensions of the components of \mathcal{O}_8 , it follows from the supergravity analysis that these conformal dimensions are also correct at strong 't Hooft coupling, though this is not guaranteed by any known non-renormalization theorem [38]. We would expect this operator to dominate the corrections to sequestering if the SUSY-breaking VEV preserves $SU(2) \times SU(2)$. Compared to the weak-coupling hidden sector expectation that dimension-two bilinears multiply visible fields and mediate soft masses $\sim |F_{hid}|^2/M_P^2$, \mathcal{O}_8 will transmit soft terms that are smaller by a factor $(\Lambda_{IR}/\Lambda_{UV})^4$.

- The lowest-dimension non-chiral supermultiplet of operators with *any* $SU(2) \times SU(2)$ quantum numbers has a lowest component of dimension $\Delta = \sqrt{28} - 2 \approx 3.29$, and a highest component of dimension $\Delta = \sqrt{28} \approx 5.29$, so we can call it $\mathcal{O}_{\sqrt{28}}$. It transforms in the $(\mathbf{3}, \mathbf{3})$ representation of $SU(2) \times SU(2)$. If the SUSY-breaking VEVs spontaneously break $SU(2) \times SU(2)$, we expect $\mathcal{O}_{\sqrt{28}}$ to mediate soft terms suppressed by a factor $(\Lambda_{IR}/\Lambda_{UV})^{1.29} \ll 1$ compared to the weak-coupling result.

Whether SUSY breaking is mediated by \mathcal{O}_8 or by $\mathcal{O}_{\sqrt{28}}$, for reasonable values of the parameters, *SUSY breaking at the end of a warped throat will be effectively sequestered from visible sector fields living on branes in the bulk of the Calabi-Yau.*

4. Gravity-side sequestering

The logic of conformal sequestering we have applied above is tight, but indirect. We will now illustrate some of the above considerations by explicit gravity-side calculation of couplings between separated D3-branes in the warped throat. We demonstrate on the gravity side the suppression of SUSY-breaking mediation to a visible-sector brane. We also discuss the effects of compactification, analogous to those of §2.2.

4.1. D3-brane couplings in a non-compact warped throat

We will now infer the coupling between a D3-brane at r_{UV} (close to r_{max}) and a probe D3-brane at the tip, $r = r_0$. To do this, we can view the D3-brane at r_{UV} as sourcing a perturbation of the throat metric (3.4), and ask how this perturbation couples to the DBI action of the IR-brane.

In addition to the metric, we will need to know the five-form flux F_5 . In the unperturbed background, in the limit where we consider the metric (3.4), this is given by

$$(F_5)_{rtx^1x^2x^3} = 4\frac{r^3}{R^4}, \quad (C_4)_{tx^1x^2x^3} = \frac{r^4}{R^4} \quad (4.1)$$

where we have chosen a gauge for C_4 . The background harmonic function

$$h(r) = \frac{R^4}{r^4} \quad (4.2)$$

is perturbed by the probe. For values of $r \ll r_{UV}$, it takes the form

$$h \rightarrow h + \delta h = \frac{R^4}{r^4} + \frac{1}{N} \frac{R^4}{r_{UV}^4} . \quad (4.3)$$

We now plug this into the Born-Infeld + Chern-Simons action of the D3-brane localized at r_0 . Because D3-branes enjoy a no-force condition in the pseudo-BPS backgrounds [18] determined by three-form fluxes in type IIB Calabi-Yau models, we find from (3.4), (4.1) that no potential is induced. However, higher derivative terms are indeed generated. The leading such operator can be understood as follows. Expanding

$$(h + \delta h)^{\pm 1/2} \approx h^{\pm 1/2} \left(1 \pm \frac{\delta h}{2h} \right) \quad (4.4)$$

we find that the modified line element in the presence of the UV brane is

$$ds^2 = h^{-1/2} \left(1 - \frac{\delta h}{2h} \right) \eta_{\mu\nu} dx^\mu dx^\nu + h^{1/2} \left(1 + \frac{\delta h}{2h} \right) \tilde{g}_{mn} dy^m dy^n . \quad (4.5)$$

This is a dilatation of the D3-brane worldvolume metric by a factor

$$\lambda = 1 - \frac{\delta h}{2h} . \quad (4.6)$$

By the AdS/CFT dictionary, this perturbation of the gravitational background should couple to a specific operator in the D3-brane field theory. For a D3-brane in an AdS throat,

the leading term in the Born-Infeld action that breaks conformal invariance and hence can couple to (4.6) is [40,41]

$$\int d^4x \left(\frac{\delta h}{2h} \right) \tilde{\mathcal{O}} (2\pi\alpha')^2 \quad (4.7)$$

on the worldvolume of the IR D3-brane, with

$$\tilde{\mathcal{O}} = \frac{2}{3} \text{tr} \left(F_{\mu\nu} F_{\rho\nu} F_{\mu\lambda} F_{\rho\lambda} + \frac{1}{2} F_{\mu\nu} F_{\rho\nu} F_{\rho\lambda} F_{\mu\lambda} - \frac{1}{4} F_{\mu\nu} F_{\mu\nu} F_{\rho\lambda} F_{\rho\lambda} - \frac{1}{8} F_{\mu\nu} F_{\rho\lambda} F_{\mu\nu} F_{\rho\lambda} \right). \quad (4.8)$$

The operator $\tilde{\mathcal{O}}$ is a component of the supermultiplet of operators

$$\mathcal{O}_8 = W_\alpha^2 \overline{W}_{\dot{\alpha}}^2 \quad (4.9)$$

which, as we have already noted [38], is the lowest-dimension supermultiplet that is invariant under $SU(2) \times SU(2)$. The fact that (4.7) is the leading coupling, and is indeed present, was tested directly by “scattering experiments” off a D3-brane in [40].

Up to the compactification effects we discuss later, this result illustrates warped sequestering at its simplest. The leading operator communicating between the two D3-branes is considerably suppressed compared to the expected quartic cross-coupling between fields on different branes. Of course, this does not demonstrate SUSY breaking, to which we now turn.

4.2. SUSY breaking and sequestering in a non-compact warped throat

A simple way of breaking SUSY in the IR of the throat is to replace the IR D3-brane above with an anti-D3-brane. The Born-Infeld logic goes through unchanged, and the leading cross-coupling is still mediated by an operator of the form (4.9). However, in this case we can borrow direct results for the probe D3/anti-D3 potential in a warped throat, as derived in [25]. The result, in the absence of a stabilizing potential for the UV D3-brane, is an attractive potential

$$V(r_{UV}) = 2T_3 \left(\frac{r_0}{R} \right)^4 \left(1 - \frac{1}{N} \frac{r_0^4}{r_{UV}^4} \right). \quad (4.10)$$

This is in agreement with our expectation that the leading operator transmitting SUSY breaking is \mathcal{O}_8 . In the presence of warping, even if the cross-coupling involved a hidden sector operator whose lowest component had dimension two (the free-field value for a chiral-antichiral bilinear), we would expect suppression of the interaction by four powers of the

IR warp-factor r_0/R , because this factor already suppresses the hidden SUSY-breaking vacuum energy density (the leading term of (4.10)) due to warping as in [42].

In (4.10), we instead see a suppression of cross-couplings by *eight* powers of the IR warp factor. This indicates that the leading hidden operator that communicates with the UV-brane fields has dimension eight (or its lowest component has dimension six). This is completely consistent with our AdS/CFT reasoning, and with the results of [41].

The fact that \mathcal{O}_8 and not $\mathcal{O}_{\sqrt{28}}$ has appeared here can be traced to our consideration of the *radial* potential. The radial position of a D3-brane corresponds to a scalar field that is invariant under the $SU(2) \times SU(2)$ global symmetry, so an operator coupling pairs of such fields, for a separated brane-antibrane pair, will also be invariant; and \mathcal{O}_8 is the lowest-dimension $SU(2) \times SU(2)$ -invariant non-chiral operator. We expect that in more general circumstances (and perhaps after including leading α' corrections to the tree-level gravity solution), couplings mediated by $\mathcal{O}_{\sqrt{28}}$ will be present.

Note that the visible brane fluctuations, corresponding to fluctuations of the UV brane position, r_{UV} , not only get highly-suppressed SUSY breaking masses by Taylor expanding (4.10) in δr_{UV} to quadratic order, but also a suppressed tadpole term linear in δr_{UV} . This reflects an instability of the visible D3-brane: the D3-brane is attracted toward the anti-D3-brane, so in principle we do not have a legitimate ground state. This unwanted feature can be cured by more realistic visible sector model-building in the Calabi-Yau bulk, and in particular by stabilization of the branes on which the visible sector is localized. Here we are only interested in illustrating the high degree of suppression of visible-sector SUSY-breaking masses.

In the earlier work [25], the focus was on using the inter-brane force (4.10) to drive inflation. The conclusion was that although the cross-couplings (4.10) are small, the full potential for r_{UV} receives additional contributions and is generically much steeper than (4.10). This could be attributed to a conformal coupling of the D3-brane modes to the four-dimensional curvature, which is nontrivial during inflation. A crucial difference is that for sequestering, one is interested in the inter-brane interactions in a background with approximately vanishing vacuum energy. Therefore, after tuning to solve the cosmological constant problem (for example, by an appropriate shift of the flux superpotential), the conformal coupling becomes a negligible effect.

4.3. Compactification effects

As in the cases of low warping in §2, we expect that compactification will result in mixing between brane fields and bulk moduli, and that generically this mixing will lead to further cross-couplings between separated brane stacks (or brane and antibrane stacks). However, locality now implies that such induced couplings must still proceed via throat modes to communicate with IR-branes or antibranes. The most important modes are still the ones we discussed above, and therefore we still expect the same degree of sequestering, although the numerical coefficients may differ. Most throat modes lead to strong sequestering of cross-couplings between UV-branes and IR-branes, and the dominant effects will arise through exchange of the Kaluza-Klein gauge bosons corresponding to the $SU(2) \times SU(2)$ isometries of the throat, which have vanishing AdS_5 mass. These are dual to the conserved currents discussed in §3.2, and thus they mediate cross-couplings in an unsuppressed fashion. However, these effects are harmless, since by the result of [16] they do not mediate IR SUSY breaking.

It is instructive to test these results directly on the gravity side. In a general warped compactification, the full Kähler potential can be rather complicated, and explicit formulae have proved elusive (see [43] for a recent discussion). However, there is a very simple system that is rich enough to manifest the desired phenomenon, the emergence of unsuppressed but nevertheless harmless cross-couplings after compactification, and for which we know the relevant formulae *exactly*. This system is the $\mathcal{N} = 4$ supersymmetric toroidal orientifold with a distribution of D3-branes discussed in [7] and in the introduction. If we separate the D3-branes into two stacks, one to provide the visible sector and the remainder to form the hidden sector, the gravitational backreaction of the latter stack forms a warped throat, as noted originally in [44]. This throat is dual to the $\mathcal{N} = 4$ SYM dynamics of the hidden stack. But despite the strong warping, the high degree of supersymmetry protects the form of the Kähler potential (1.4), and it is unaltered by the continuation to large $g_s N$. As noted in the introduction, this system does have formally unsequestered cross-couplings. But, these are precisely couplings of the visible sector to the conserved currents of the hidden sector, associated to the throat isometries, and in the presence of hidden SUSY breaking these couplings pose no threat (at least in the circumstances described in [16]). The fact that even in the unwarped regime the $\mathcal{N} = 4$ orientifold enjoys this milder form of sequestering is “explained” by the fact that it *must* do so in the warped regime $g_s N \gg 1$ (by our previous arguments), and the non-renormalization of the Kähler

potential then forces it to do so even in the unwarped limit. One can also directly see that the cross-couplings vanish in the non-compact limit (in which the radii are taken to infinity), which agrees with the observation of §2.1 that the D3-branes do not see each other prior to compactification.

5. Classes of warped sequestering

In our discussion of warped sequestering, we have focused on concrete examples arising in compactifications that incorporate the warped deformed conifold [17]. However, while this is a useful and common example, we wish to emphasize that there is by now a wide spectrum of other throats (understood in different levels of detail) that could be more useful for various model building purposes.

The two main ingredients that go into specifying a model, independent of visible-sector model building, are: (a) the description of an appropriate AdS/CFT dual pair that governs the throat dynamics (describes the approximate fixed point controlling the RG cascade), and (b) the incorporation of an appropriate SUSY-breaking sector in the IR region of the throat. Both are subjects of intense research, characterized by rapid recent progress.

In considering possibilities for (a), it is useful to keep in mind that the main dangers to sequestering involve global currents in the CFT. Any $\mathcal{N} = 1$ CFT has a $U(1)_R$ symmetry, which does not pose any danger to sequestering. But one could wish to find examples that lack additional continuous global symmetries, because when such symmetries are present we must take extra care to ensure that deformations that break hidden sector conformal invariance are also symmetric [16] (as we checked in the case of the conifold throat). The smaller the global symmetry group, the fewer such checks we need to perform. Several new classes of Sasaki-Einstein solutions to string theory have been discovered recently, including the $Y^{p,q}$ and $L^{p,q,r}$ series [45,46]. Similarly, cascades based on diverse solutions have been described in *e.g.* [47,48,49,50]. While the relevant CFTs in these new classes are characterized by less symmetry than the conifold, they still possess continuous global symmetries beyond the minimal $U(1)_R$.

However, it is possible to exhibit simple, infinite families of dual pairs that only possess $U(1)_R$. Even $\mathcal{N} = 1$ orbifolds of the $\mathcal{N} = 4$ SYM theory [51,52,53] by non-Abelian discrete subgroups of $SU(3)$ are useful in this regard. Indeed, the series of discrete groups Δ_{3n^2} and Δ_{6n^2} give rise to orbifold CFTs with only the $U(1)_R$ subgroup of the $SO(6)$ symmetry

of $\mathcal{N} = 4$ SYM surviving (the relevant quiver gauge theories were derived in detail in [54]). One can see this as follows. The group Δ_{3n^2} , for instance, is generated by

$$\alpha : (z_1, z_2, z_3) \rightarrow (e^{2\pi i/n} z_1, e^{-2\pi i/n} z_2, z_3) , \quad (5.1)$$

$$\beta : (z_1, z_2, z_3) \rightarrow (z_1, e^{2\pi i/n} z_2, e^{-2\pi i/n} z_3) , \quad (5.2)$$

and

$$\gamma : (z_1, z_2, z_3) \rightarrow (z_3, z_1, z_2) . \quad (5.3)$$

The maximal abelian subgroup of $SO(6)$ is generated by independent phase rotations of the z_i . It is easy to see that γ only commutes with the single $U(1)$ that rotates the z_i by equal phases. Therefore, these nonabelian orbifolds preserve only a single $U(1)_R$ subgroup of the $SO(6)$ symmetry of $\mathcal{N} = 4$ SYM. It would be interesting to make cascading solutions controlled by these CFTs but ending smoothly in the IR [55]. These series in some sense provide a counterexample to the expectation in [37] that typical CFTs will contain additional continuous symmetries, with the corresponding dangerous supermultiplets of currents.

Supersymmetry breaking at the end of RG cascades (or at the end of warped throats, in the dual gravity language) has also been a focus of recent research. In the warped deformed conifold, the models of [19] provide examples. The quivers that characterize simple \mathbb{Z}_k quotients of the conifold were described in [56], and cascades governed by the associated CFTs together with SUSY breaking mechanisms at the end of the throat have been investigated in [57]. Other cascades with SUSY breaking, associated with cones over del Pezzo surfaces, have been examined in [49,58,59]. Many of these models suffer from a runaway vacuum [60]; for the state of the art in obtaining stable vacua this way, see [61,62,63].

Finally, although the bulk of the literature on warped compactifications in string theory focuses on type IIB vacua, this is only for reasons of convenience: the solutions are conformally Calabi-Yau and the most familiar examples of AdS/CFT arise in this context. We expect that further research will uncover similar rich families of warped compactifications in the type IIA, heterotic, and eleven-dimensional limits, which could be equally or more promising for the construction of fully realistic models.

6. Conclusion

We have argued that warped sequestering is robustly attainable in string theory, and that sequestered interactions between D3-branes are possible even in some unwarped models. Warped throats are a rather generic feature of type IIB flux compactifications [17,18,64,65,66], so we conclude that one should be able to design, without undue difficulty, concrete string theory models of SUSY breaking in which the flavor problem is vitiated by sequestering. Such models merit further detailed exploration.

One of our motivations for pursuing this work was the advent of mirage mediation [4], a scenario where anomaly mediation is combined with moduli mediation in a warped flux compactification. We hope to have provided a firm foundation for discussion of the circumstances under which such models can lead to flavor-blind supersymmetry breaking. We note that the sequestering due to warped throats is insensitive to the details of the standard model construction in the bulk of the Calabi-Yau space. In particular, while D7-brane fields are not sequestered from D3-brane fields in the non-compact, unwarped limit, SUSY breaking at the end of a throat can be sequestered from D7-branes in a suitably warped type IIB flux compactification. This makes it seem likely that relatively realistic models of warped sequestering incorporating unification of coupling constants can be constructed in string theory.

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References

- [1] L. Randall and R. Sundrum, “Out of this world supersymmetry breaking,” Nucl. Phys. B **557**, 79 (1999) [arXiv:hep-th/9810155].
- [2] G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, “Gaugino mass without singlets,” JHEP **9812**, 027 (1998) [arXiv:hep-ph/9810442].
- [3] D. E. Kaplan, G. D. Kribs and M. Schmaltz, “Supersymmetry breaking through transparent extra dimensions,” Phys. Rev. D **62**, 035010 (2000) [arXiv:hep-ph/9911293]; Z. Chacko, M. A. Luty, A. E. Nelson and E. Ponton, “Gaugino mediated supersymmetry breaking,” JHEP **0001**, 003 (2000) [arXiv:hep-ph/9911323].
- [4] K. Choi, A. Falkowski, H. P. Nilles and M. Olechowski, “Soft supersymmetry breaking in KKL T flux compactification,” Nucl. Phys. B **718**, 113 (2005) [arXiv:hep-th/0503216].
- [5] A. Pomarol and R. Rattazzi, “Sparticle masses from the superconformal anomaly,” JHEP **9905**, 013 (1999) [arXiv:hep-ph/9903448]; Z. Chacko, M. A. Luty, I. Maksymyk and E. Ponton, “Realistic anomaly-mediated supersymmetry breaking,” JHEP **0004**, 001 (2000) [arXiv:hep-ph/9905390]; E. Katz, Y. Shadmi and Y. Shirman, “Heavy thresholds, slepton masses and the μ term in anomaly mediated supersymmetry breaking,” JHEP **9908**, 015 (1999) [arXiv:hep-ph/9906296]; K. I. Izawa, Y. Nomura and T. Yanagida, “Cosmological constants as messenger between branes,” Prog. Theor. Phys. **102**, 1181 (1999) [arXiv:hep-ph/9908240]; M. Carena, K. Huitu and T. Kobayashi, “RG-invariant sum rule in a generalization of anomaly mediated SUSY breaking models,” Nucl. Phys. B **592**, 164 (2001) [arXiv:hep-ph/0003187]; B. C. Allanach and A. Dedes, “R-parity violating anomaly mediated supersymmetry breaking,” JHEP **0006**, 017 (2000) [arXiv:hep-ph/0003222]; I. Jack and D. R. T. Jones, “R-symmetry, Yukawa textures and anomaly mediated supersymmetry breaking,” Phys. Lett. B **491**, 151 (2000) [arXiv:hep-ph/0006116]; D. E. Kaplan and G. D. Kribs, “Gaugino-assisted anomaly mediation,” JHEP **0009**, 048 (2000) [arXiv:hep-ph/0009195]; N. Arkani-Hamed, D. E. Kaplan, H. Murayama and Y. Nomura, “Viable ultraviolet-insensitive supersymmetry breaking,” JHEP **0102**, 041 (2001) [arXiv:hep-ph/0012103]; Z. Chacko and M. A. Luty, “Realistic anomaly mediation with bulk gauge fields,” JHEP **0205**, 047 (2002) [arXiv:hep-ph/0112172]; A. E. Nelson and N. J. Weiner, “Gauge/anomaly Syzygy and generalized brane world models of supersymmetry breaking,” Phys. Rev. Lett. **88**, 231802 (2002) [arXiv:hep-ph/0112210].
- [6] M. A. Luty and R. Sundrum, “Radius stabilization and anomaly-mediated supersymmetry breaking,” Phys. Rev. D **62**, 035008 (2000) [arXiv:hep-th/9910202].
- [7] A. Anisimov, M. Dine, M. Graesser and S. D. Thomas, “Brane world SUSY breaking,” Phys. Rev. D **65**, 105011 (2002) [arXiv:hep-th/0111235];

- A. Anisimov, M. Dine, M. Graesser and S. D. Thomas, “Brane world SUSY breaking from string/M theory,” JHEP **0203**, 036 (2002) [arXiv:hep-th/0201256].
- [8] S. Kachru, J. McGreevy and P. Svrcek, “Bounds on masses of bulk fields in string compactifications,” JHEP **0604**, 023 (2006) [arXiv:hep-th/0601111].
 - [9] M. A. Luty and R. Sundrum, “Hierarchy stabilization in warped supersymmetry,” Phys. Rev. D **64**, 065012 (2001) [arXiv:hep-th/0012158].
 - [10] J. M. Maldacena, “The large N limit of superconformal field theories and supergravity,” Adv. Theor. Math. Phys. **2**, 231 (1998) [Int. J. Theor. Phys. **38**, 1113 (1999)] [arXiv:hep-th/9711200].
 - [11] S. S. Gubser, I. R. Klebanov and A. M. Polyakov, Phys. Lett. B **428**, 105 (1998) [arXiv:hep-th/9802109];
E. Witten, “Anti-de Sitter space and holography,” Adv. Theor. Math. Phys. **2**, 253 (1998) [arXiv:hep-th/9802150].
 - [12] M. A. Luty and R. Sundrum, “Supersymmetry breaking and composite extra dimensions,” Phys. Rev. D **65**, 066004 (2002) [arXiv:hep-th/0105137].
 - [13] A. E. Nelson and M. J. Strassler, “Exact results for supersymmetric renormalization and the supersymmetric flavor problem,” JHEP **0207**, 021 (2002) [arXiv:hep-ph/0104051];
A. E. Nelson and M. J. Strassler, “Suppressing flavor anarchy,” JHEP **0009**, 030 (2000) [arXiv:hep-ph/0006251].
 - [14] M. Luty and R. Sundrum, “Anomaly mediated supersymmetry breaking in four dimensions, naturally,” Phys. Rev. D **67**, 045007 (2003) [arXiv:hep-th/0111231].
 - [15] M. Ibe, K. Izawa, Y. Nakayama, Y. Shinbara and T. Yanagida, “Conformally sequestered SUSY breaking in vector-like gauge theories,” Phys. Rev. D **73**, 015004 (2006) [arXiv:hep-ph/0506023];
M. Ibe, K. Izawa, Y. Nakayama, Y. Shinbara and T. Yanagida, “More on conformally sequestered SUSY breaking,” Phys. Rev. D **73**, 035012 (2006) [arXiv:hep-ph/0509229].
 - [16] M. Schmaltz and R. Sundrum, “Conformal sequestering simplified,” JHEP **0611**, 011 (2006) [arXiv:hep-th/0608051].
 - [17] I. R. Klebanov and M. J. Strassler, “Supergravity and a confining gauge theory: Duality cascades and chiSB-resolution of naked singularities,” JHEP **0008**, 052 (2000) [arXiv:hep-th/0007191].
 - [18] S. B. Giddings, S. Kachru and J. Polchinski, “Hierarchies from fluxes in string compactifications,” Phys. Rev. D **66**, 106006 (2002) [arXiv:hep-th/0105097].
 - [19] S. Kachru, J. Pearson and H. L. Verlinde, “Brane/flux annihilation and the string dual of a non-supersymmetric field theory,” JHEP **0206**, 021 (2002) [arXiv:hep-th/0112197].
 - [20] K. Becker and M. Becker, “M-Theory on Eight-Manifolds,” Nucl. Phys. B **477**, 155 (1996) [arXiv:hep-th/9605053].

- [21] S. Gukov, C. Vafa and E. Witten, “CFT’s from Calabi-Yau four-folds,” Nucl. Phys. B **584**, 69 (2000) [Erratum-ibid. B **608**, 477 (2001)] [arXiv:hep-th/9906070].
- [22] K. Dasgupta, G. Rajesh and S. Sethi, “M theory, orientifolds and G-flux,” JHEP **9908**, 023 (1999) [arXiv:hep-th/9908088].
- [23] M. Grana, T. W. Grimm, H. Jockers and J. Louis, “Soft supersymmetry breaking in Calabi-Yau orientifolds with D-branes and fluxes,” Nucl. Phys. B **690**, 21 (2004) [arXiv:hep-th/0312232].
- [24] O. DeWolfe and S. B. Giddings, “Scales and hierarchies in warped compactifications and brane worlds,” Phys. Rev. D **67**, 066008 (2003) [arXiv:hep-th/0208123].
- [25] S. Kachru, R. Kallosh, A. Linde, J. M. Maldacena, L. McAllister and S. P. Trivedi, “Towards inflation in string theory,” JCAP **0310**, 013 (2003) [arXiv:hep-th/0308055].
- [26] H. Jockers and J. Louis, “The effective action of D7-branes in $N = 1$ Calabi-Yau orientifolds,” Nucl. Phys. B **705**, 167 (2005) [arXiv:hep-th/0409098].
- [27] S. B. Giddings and A. Maharana, “Dynamics of warped compactifications and the shape of the warped landscape,” Phys. Rev. D **73**, 126003 (2006) [arXiv:hep-th/0507158].
- [28] D. Baumann, A. Dymarsky, I. R. Klebanov, J. Maldacena, L. McAllister and A. Murugan, “On D3-brane potentials in compactifications with fluxes and wrapped D-branes,” JHEP **0611**, 031 (2006) [arXiv:hep-th/0607050].
- [29] O. J. Ganor, “A note on zeroes of superpotentials in F-theory,” Nucl. Phys. B **499**, 55 (1997) [arXiv:hep-th/9612077].
- [30] P. Candelas, A. Font, S. H. Katz and D. R. Morrison, “Mirror symmetry for two parameter models. 2,” Nucl. Phys. B **429**, 626 (1994) [arXiv:hep-th/9403187];
F. Denef, M. R. Douglas and B. Florea, “Building a better racetrack,” JHEP **0406**, 034 (2004) [arXiv:hep-th/0404257];
V. Balasubramanian, P. Berglund, J. P. Conlon and F. Quevedo, “Systematics of moduli stabilisation in Calabi-Yau flux compactifications,” JHEP **0503**, 007 (2005) [arXiv:hep-th/0502058].
- [31] K. Choi and K. S. Jeong, “Supersymmetry breaking and moduli stabilization with anomalous $U(1)$ gauge symmetry,” JHEP **0608**, 007 (2006) [arXiv:hep-th/0605108].
- [32] M. Berg, M. Haack and B. Kors, “String loop corrections to Kaehler potentials in orientifolds,” JHEP **0511**, 030 (2005) [arXiv:hep-th/0508043];
M. Berg, M. Haack and B. Kors, “On volume stabilization by quantum corrections,” Phys. Rev. Lett. **96**, 021601 (2006) [arXiv:hep-th/0508171].
- [33] R. Blumenhagen, M. Cvetič, P. Langacker and G. Shiu, “Toward realistic intersecting D-brane models,” Ann. Rev. Nucl. Part. Sci. **55**, 71 (2005) [arXiv:hep-th/0502005].
- [34] F. Marchesano, “Progress in D-brane model building,” arXiv:hep-th/0702094.
- [35] I. R. Klebanov and E. Witten, “Superconformal field theory on threebranes at a Calabi-Yau singularity,” Nucl. Phys. B **536**, 199 (1998) [arXiv:hep-th/9807080].

- [36] S. Kachru, R. Kallosh, A. Linde and S. P. Trivedi, “de Sitter vacua in string theory,” *Phys. Rev. D* **68**, 046005 (2003) [arXiv:hep-th/0301240].
- [37] M. Dine, P. J. Fox, E. Gorbatov, Y. Shadmi, Y. Shirman and S. D. Thomas, “Visible effects of the hidden sector,” *Phys. Rev. D* **70**, 045023 (2004) [arXiv:hep-ph/0405159].
- [38] A. Ceresole, G. Dall’Agata, R. D’Auria and S. Ferrara, “Spectrum of type IIB supergravity on $\text{AdS}(5) \times \text{T}(11)$: Predictions on $N = 1$ SCFT’s,” *Phys. Rev. D* **61**, 066001 (2000) [arXiv:hep-th/9905226];
A. Ceresole, G. Dall’Agata and R. D’Auria, “KK spectroscopy of type IIB supergravity on $\text{AdS}(5) \times \text{T}(11)$,” *JHEP* **9911**, 009 (1999) [arXiv:hep-th/9907216].
- [39] O. Aharony, Y. E. Antebi and M. Berkooz, “Open string moduli in KKLt compactifications,” *Phys. Rev. D* **72**, 106009 (2005) [arXiv:hep-th/0508080].
- [40] S. S. Gubser, A. Hashimoto, I. R. Klebanov and M. Krasnitz, “Scalar absorption and the breaking of the world volume conformal invariance,” *Nucl. Phys. B* **526**, 393 (1998) [arXiv:hep-th/9803023].
- [41] S. Ferrara, M. A. Lledo and A. Zaffaroni, “Born-Infeld corrections to D3 brane action in $\text{AdS}(5) \times \text{S}(5)$ and $N = 4$, $d = 4$ primary superfields,” *Phys. Rev. D* **58**, 105029 (1998) [arXiv:hep-th/9805082].
- [42] L. Randall and R. Sundrum, “A large mass hierarchy from a small extra dimension,” *Phys. Rev. Lett.* **83**, 3370 (1999) [arXiv:hep-ph/9905221].
- [43] C. P. Burgess, P. G. Camara, S. P. de Alwis, S. B. Giddings, A. Maharana, F. Quevedo and K. Suruliz, “Warped supersymmetry breaking,” arXiv:hep-th/0610255.
- [44] H. L. Verlinde, “Holography and compactification,” *Nucl. Phys. B* **580**, 264 (2000) [arXiv:hep-th/9906182].
- [45] J. P. Gauntlett, D. Martelli, J. Sparks and D. Waldram, “Sasaki-Einstein metrics on $\text{S}(2) \times \text{S}(3)$,” *Adv. Theor. Math. Phys.* **8**, 711 (2004) [arXiv:hep-th/0403002].
- [46] M. Cvetič, H. Lu, D. N. Page and C. N. Pope, “New Einstein-Sasaki spaces in five and higher dimensions,” *Phys. Rev. Lett.* **95**, 071101 (2005) [arXiv:hep-th/0504225].
- [47] C. P. Herzog, Q. J. Ejaz and I. R. Klebanov, “Cascading RG flows from new Sasaki-Einstein manifolds,” *JHEP* **0502**, 009 (2005) [arXiv:hep-th/0412193].
- [48] S. Franco, A. Hanany and A. M. Uranga, “Multi-flux warped throats and cascading gauge theories,” *JHEP* **0509**, 028 (2005) [arXiv:hep-th/0502113].
- [49] M. Bertolini, F. Bigazzi and A. L. Cotrone, “Supersymmetry breaking at the end of a cascade of Seiberg dualities,” *Phys. Rev. D* **72**, 061902 (2005) [arXiv:hep-th/0505055].
- [50] C. Doran, M. Headrick, C. P. Herzog, J. Kantor and T. Wiseman, “Numerical Kaehler-Einstein metric on the third del Pezzo,” arXiv:hep-th/0703057.
- [51] M. R. Douglas, B. R. Greene and D. R. Morrison, “Orbifold resolution by D-branes,” *Nucl. Phys. B* **506**, 84 (1997) [arXiv:hep-th/9704151].
- [52] S. Kachru and E. Silverstein, “4d conformal theories and strings on orbifolds,” *Phys. Rev. Lett.* **80**, 4855 (1998) [arXiv:hep-th/9802183].

- [53] A. E. Lawrence, N. Nekrasov and C. Vafa, “On conformal field theories in four dimensions,” Nucl. Phys. B **533**, 199 (1998) [arXiv:hep-th/9803015].
- [54] B. R. Greene, C. I. Lazaroiu and M. Raugas, “D-branes on nonabelian threefold quotient singularities,” Nucl. Phys. B **553**, 711 (1999) [arXiv:hep-th/9811201].
- [55] W. Chuang and M. Mulligan, work in progress.
- [56] A. M. Uranga, “Brane configurations for branes at conifolds,” JHEP **9901**, 022 (1999) [arXiv:hep-th/9811004].
- [57] R. Argurio, M. Bertolini, S. Franco and S. Kachru, “Gauge/gravity duality and metastable dynamical supersymmetry breaking,” arXiv:hep-th/0610212; and to appear.
- [58] D. Berenstein, C. P. Herzog, P. Ouyang and S. Pinansky, “Supersymmetry breaking from a Calabi-Yau singularity,” JHEP **0509**, 084 (2005) [arXiv:hep-th/0505029].
- [59] S. Franco, A. Hanany, F. Saad and A. M. Uranga, “Fractional branes and dynamical supersymmetry breaking,” JHEP **0601**, 011 (2006) [arXiv:hep-th/0505040].
- [60] K. Intriligator and N. Seiberg, “The runaway quiver,” JHEP **0602**, 031 (2006) [arXiv:hep-th/0512347].
- [61] R. Argurio, M. Bertolini, C. Closset and S. Cremonesi, “On stable non-supersymmetric vacua at the bottom of cascading theories,” JHEP **0609**, 030 (2006) [arXiv:hep-th/0606175].
- [62] D. E. Diaconescu, R. Donagi and B. Florea, “Metastable quivers in string compactifications,” arXiv:hep-th/0701104.
- [63] M. Wijnholt, “Geometry of Particle Physics,” arXiv:hep-th/0703047.
- [64] F. Denef and M. R. Douglas, “Distributions of flux vacua,” JHEP **0405**, 072 (2004) [arXiv:hep-th/0404116].
- [65] A. Giryavets, S. Kachru and P. K. Tripathy, “On the taxonomy of flux vacua,” JHEP **0408**, 002 (2004) [arXiv:hep-th/0404243].
- [66] A. Hebecker and J. March-Russell, “The ubiquitous throat,” arXiv:hep-th/0607120.